Gamma ray bursts from superconducting cosmic strings

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Abstract

Cusps of superconducting strings can serve as GRB engines. A powerful beamed pulse of electromagnetic radiation from a cusp produces a jet of accelerated particles, whose propagation is terminated by the shock responsible for GRB. A single free parameter, the string scale of symmetry breaking $\eta \sim 10^{14}\ GeV$, together with reasonable assumptions about the magnitude of cosmic magnetic fields and the fraction of volume that they occupy, explains the GRB rate, duration and fluence, as well as the observed ranges of these quantities. The wiggles on the string can drive the short-time structures of GRB. This model predicts that GRBs are accompanied by strong bursts of gravitational radiation which should be detectable by LIGO, VIRGO and LISA detectors. Another prediction is the diffuse X- and gamma-ray radiation at 8 MeV - 100 GeV with a spectrum and flux comparable to the observed. The weakness of the model is the prediction of too low rate of GRBs from galaxies, as compared with observations. This suggests that either the capture rate of string loops by galaxies is underestimated in our model, or that GRBs from cusps are responsible for only a subset of the observed GRBs not

associated with galaxies.

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I. INTRODUCTION

Existing models of gamma ray bursts (GRBs) face the problem of explaining the tremendous energy released by the central engine [1]. In the case of isotropic emission, the total energy output should be as high as 1.4×10^{54} ergs, in case of GRB 990123 with redshift $z = 1.6$. Strongly beamed emission is needed for all known engine models, such as mergers and hypernovae, but such extreme beaming is difficult to arrange (see recent discussion by Blandford [2] and Rees [3]). In this paper we show that emission of pulsed electromagnetic radiation from cusps of superconducting cosmic strings naturally solves this problem and explains the observational GRB data using only one engine parameter [4].

Cosmic strings are linear defects that could be formed at a symmetry breaking phase transition in the early universe [5]. Strings predicted in most grand unified models respond to external electromagnetic fields as thin superconducting wires [6]. As they move through cosmic magnetic fields, such strings develop electric currents. Oscillating loops of superconducting string emit short bursts of highly beamed electromagnetic radiation through small string segments, centered at peculiar points on a string, cusps, where velocity reaches speed of light. [7,8].

The idea that GRBs could be produced at cusps of superconducting strings was first suggested by Babul, Paczynski and Spergel [9] (BPS) and further explored by Paczynski [10]. They assumed that the bursts originate at very high redshifts ($z \sim 100 - 1000$), with GRB photons produced either directly or in electromagnetic cascades developing due to interaction with the microwave background. This model requires the existence of a strong primordial magnetic field to generate the string currents.

As it stands, the BPS model does not agree with observations. The observed GRB redshifts are in the range $z \lesssim 3$, and the observed duration of the bursts ($10^{-2}s \lesssim \tau \lesssim 10^{3}s$) is significantly longer than that predicted by the model. On the theoretical side, our understanding of cosmic string evolution and of the GRB generation in relativistic jets have

considerably evolved since the BPS papers were written. Our goal in this paper is to revive the BPS idea taking stock of these recent advances.

As in the BPS model we shall use the cusp of a superconducting string as the central engine in GRB. It provides the tremendous engine energy naturally beamed. Our main observation is that putting superconducting cusps in a different enviroment, the magnetized plasma at a relatively small redshift z, results in a different mechanism of gamma radiation, which leads to a good agreement with GRB observational data.

GRB radiation in our model arises as follows. Low-frequency electromagnetic radiation from a cusp loses its energy by accelerating particles of the plasma to very large Lorentz factors. Like the initial electromagnetic pulse, the particles are beamed and give rise to a hydrodynamical flow in the surrounding gas, terminated by a shock, as in the standard fireball theory of GRB [11] (for a review see [1]).

The string symmetry breaking scale η will be the only string parameter used in our calculations. With reasonable assumptions about the magnitude of cosmic magnetic fields and the fraction of volume in the Universe that they occupy, this parameter is sufficient to account for all main GRB observational quantities: the duration τ_{GRB}, the rate of events $\dot{N}_{GRB}$, and the fluence S.

We begin in the next Section with a brief review of cosmic string properties and evolution, with an emphasis on the physics of cusps and on the generation and dissipation of electric current in superconducting strings. (The discussion of the latter topic in the existing literature is often oversimplified and sometimes incorrect, so we review it in more detail than we otherwise would.) The GRB characteristics in our model are calculated in Section III, and the hydrodynamical aspects of the model are discussed in Section IV. In Section V we discuss the diffuse X-ray and γ-ray backgrounds predicted by the model, as well as other observational predictions, which include GRB repeaters, bursts of gravitational radiation, and ultrahigh-energy particles.

II. STRING OVERVIEW

A. String properties and evolution

Here we briefly review some aspects of cosmic string properties and evolution, which are relevant for the discussion below (for a detailed review and references see [5]).

Strings are characterized by the energy scale of symmetry breaking η, which is given by the expectation value of the corresponding Higgs field, $\langle\phi\rangle = \eta$. The mass per unit length of string is given by

$$\mu \sim \eta^2. \tag{1}$$

An important dimensionless parameter characterizing the gravitational interactions of strings is

$$G\mu \sim (\eta/m_P)^2, \tag{2}$$

where G is Newton's constant and m_P is the Planck mass. In many models this is the only relevant string parameter.

Numerical simulations of cosmic string evolution indicate that strings evolve in a self-similar manner [12–14]. A horizon-size volume at any time t contains a few long strings stretching across the volume and a large number of small closed loops. The typical distance between long strings and their characteristic curvature radius are both $\sim t$, but, in addition, the strings have small-scale wiggles of wavelength down to

$$l \sim \alpha t, \tag{3}$$

with $\alpha \ll 1$. The typical length of loops being chopped off the long strings is comparable to the scale of the smallest wiggles (3).

The loops oscillate periodically and lose their energy, mostly by gravitational radiation. For a loop of invariant length l [15], the oscillation period is $T_l = l/2$ and the lifetime is $\tau_l \sim l/k_g G\mu$, where $k_g \sim 50$ is a numerical coefficient.

The exact value of the parameter α in (3) is not known. Numerical simulations give only an upper bound, $\alpha \lesssim 10^{-3}$, while the analysis of gravitational radiation backreaction indicates that $\alpha \gtrsim k_g G\mu$. We shall assume, following [12], that α is determined by the gravitational backreaction, so that

$$\alpha \sim k_g G\mu. \tag{4}$$

Note that in this case the loops decay within about one Hubble time of their formation. Then, most of the loops at time t have length $l \sim \alpha t$, and their number density is given by

$$n_l(t) \sim \alpha^{-1} t^{-3}. \tag{5}$$

Analysis of string equations of motion reveals that oscillating loops tend to form cusps, where for a brief period of time the string reaches the speed very close to the speed of light. Near a cusp, the string gets contracted by a large factor, its rest energy being turned into kinetic energy. For a string segment of invariant length $\delta l \ll l$, the maximum contraction factor is $\sim l/\delta l$, resulting in a Lorentz factor

$$\gamma \sim l/\delta l. \tag{6}$$

To avoid confusion, we note that cusps were originally defined [16] as points of infinite contraction, where the string momentarily reaches the speed of light. Strictly speaking, such cusps can be formed only on idealized infinitely thin strings. For realistic strings, the cusp development is truncated either by the annihilation of overlapping string segments at the tip of the cusp [17–19] or for superconducting strings, by the back reaction of charge carriers or of the electromagnetic radiation. However, unless the string current is very large, so that the energy of the charge carriers is comparable to that of the string itself, the truncation occurs at a very large Lorentz factor and the string exhibits cusp-like behavior. Below we shall use the word "cusps" to refer to such ultra-relativistic string segments.

Cusps typically form a few times during an oscillation period, but it is possible to construct (somewhat contrived) loop configurations exhibiting no cusps. Apart from various backreaction effects, the motion of loops is strictly periodic, and thus cusps reappear at nearly the same locations on the string in each oscillation period.

Another peculiar feature that one can expect to find on string loops is a kink [20]. It is characterized by a sharp bent, where the string direction changes discontinuously. Two oppositely moving kinks are produced on a loop at the moment when the loop is disconnected from a long string. The kinks then run around the loop at the speed of light.

B. String superconductivity

As first shown by Witten [6], strings predicted in a wide class of elementary particle models behave as superconducting wires. If some fermions acquire their mass as a result of the same symmetry breaking that is responsible for the string formation, then these fermions are massive outside the string but are massless inside. If in addition some of these fermions are electrically charged, then the strings have massless charge carriers which travel along the string at the speed of light. The fermion mass outside the string is $m = \lambda\eta$, where λ is the Yukawa coupling of the fermion to the Higgs field of the string. Yukawa couplings in particle physics models are often very small, so it is not unusual to have $m \ll \eta$. String superconductivity can also be bosonic, with charge carriers being either spin-zero bosons or spin-one gauge particles. Here, we shall consider only fermionic superconductivity.

An electric field E applied along a superconducting string generates an electric current. The Fermi momentum of the charge carriers grows with time as $\dot{p}_F = eE$, where e is the elementary charge, and the number of fermions per unit length, $n = p_F/2\pi$, also grows, $\dot{n} \sim eE$. The resulting current, $J \sim en$, grows at the rate

$$dJ/dt \sim e^2 E. \tag{7}$$

A superconducting loop oscillating in a magnetic field B acts as an *ac* generator and develops an *ac* current of amplitude

$$J_0 \sim e^2 B l. \tag{8}$$

This loop current is not homogeneous; it changes direction along the string and is more accurately described as current-charge oscillations. Some portions of the loop develop charge densities $\sim J_0$. For typical values used in the calculations below, $B = 1 \cdot 10^{-7}$ G and $l = \alpha t_0 \sim 30$ pc, with $\alpha = 1 \cdot 10^{-8}$ and $t_0 \sim 10^{10}$ yrs the present age of the universe, one obtains $J_0 \sim 2 \cdot 10^5$ GeV.

The local value of the string current can be greatly enhanced in the vicinity of cusps. The portion of the string that attains a Lorentz factor γ is contracted by a factor $\sim 1/\gamma$. The charge carrier density, and thus the current, are enhanced by the same factor, so the current becomes (in the local rest frame of the string)

$$J_\gamma \sim \gamma J_0. \tag{9}$$

The growth of electric current at the cusp is terminated at a critical value J_{max} when the energy of charge carriers becomes comparable to that of the string itself, $(J/e)^2 \sim \mu$. This gives J_{max} and γ_{max} as [21]

$$J_{max} \sim e\eta, \qquad \gamma_{max} \sim (e\eta/J_0). \tag{10}$$

Alternatively, the cusp development can be terminated by small-scale wiggles on the string [22]. If the wiggles contribute a fraction $\epsilon \ll 1$ to the total energy of the string, then the maximum Lorentz factor is less than (10), and is given by

$$\gamma_{max} \sim \epsilon^{-1/2}. \tag{11}$$

The actual value of γ_{max} is not important for most of the following discussion.

In realistic models, the strings have several fermion species as charge carriers. It can be shown that fermions of a given species can move only in a certain direction along the string. Thus, the charge carrier species can be divided into left-movers and right-movers. If, for example, the applied electric field is directed to the right, it produces positively charged right-movers and negatively charged left-movers (and vice versa for the opposite direction of E). The left-movers and right-movers usually differ by flavour, lepton number, or some other conserved or weakly violated quantum number.

In the absence of an external electromagnetic field, the current in an oscillating loop decays due to various dissipation mechanisms. These include: scattering of left- and right-movers [23,24], electromagnetic back-reaction [25,26], and plasma effects [27].

Charge carrier loss due to scattering of left- and right-movers is highly model-dependent. If the scattering is mediated by superheavy gauge bosons of mass $M_X \sim 10^{15}$ GeV, then the characteristic scattering time is [23]

$$\tau_{sc} \sim 3 \cdot 10^4 \left(\frac{J}{10^5 \text{ GeV}} \right)^{-5} \text{ yrs.} \tag{12}$$

For $J < 10^4$ GeV this time is greater than the age of the universe, but τ_{sc} decreases rapidly with the growth of the current and becomes comparable to the typical oscillation period of loops ($T_l \sim 100$ yrs for $l \sim 30$ pc) for $J \sim 3 \cdot 10^5$ GeV. In near-cusp regions, where $J \gg 10^5$ GeV, charge carrier scattering becomes very efficient.

We note, however, that this current loss mechanism has an important limitation. The densities of left- and right-moving charge carriers are typically not equal, and even if scattering were 100% efficient, it would stop after eliminating the minority charge carriers, leaving the string with a chiral current (that is, with a current consisting of only left- or right-movers). This is what we expect to happen in the vicinity of cusps.

The electromagnetic back-reaction typically damps the loop current on a timescale $\tau_{em} \sim l/e^2 \sim 100l$, which is much shorter than the loop's lifetime. It tends to damp the spatial component of the current, with the total charge of the loop remaining the same, so in

the absence of other effects the end result would be left- and right-moving currents of the same magnitude and the same charge.[1] Combined with scattering of charge carriers, this mechanism can dissipate loop charges and currents, even in the chiral case. Moreover, the string charge is almost completely screened by a vacuum condensate [28], so the string is effectively neutral even if the scattering rate is low and there is some residual charge. It should be noted that the physics of electromagnetic back-reaction can be significantly modified by plasma effects, which are presently not well understood. Thompson [27] has argued that current damping becomes more efficient in the presence of plasma.

Another mechanism that can dissipate a large chiral current operates when a loop oscillates in an external magnetic field. The emf induced in the loop oscillates with the same period. Suppose for definiteness that the loop initially has a chiral current J_i consisting of positively charged right-movers. When the emf is directed oppositely to this current, the magnitude of the right-moving current is reduced by $\sim J_0$ and a positively charged left-moving current of magnitude $\sim J_0$ is generated, with J_0 from Eq. (8). Left- and right-movers can now scatter off the string, and if $\tau_{sc} < T_l$, the chiral component of the current will be reduced by $\sim J_0$. The initial current will then be dissipated in $\sim J_i/J_0$ oscillations.

The effect of all these dissipation mechanisms is to damp the loop's charge and current on a timescale

$$\tau_d \sim (1 - 100)l. \tag{13}$$

This means in particular that the loop quickly forgets any initial charge or current that it inherits when it is chopped off the long string network. The magnitude of the current in a

[1]Spergel *et. al.* [26] argued that the *dc* component of the current cannot be changed by the e-m back-reaction. However, their Eq. (11) which they quote in support of this statement is in fact an expression of charge conservation.

loop is determined mainly by the local magnitude of the cosmic magnetic field, as in Eq. (8).

We note finally that Eq. (8) for the current is modified when the loop has an appreciable center-of-mass velocity v [8]. In this case, the amplitude of current-charge oscillations grows linearly with time, until the growth is hampered by the damping processes. The resulting amplitude is

$$J_0 \sim e^2 B v \tau_d. \tag{14}$$

Loops can have high center-of-mass velocities, $v \sim 1$, but in view of the uncertainty in the damping time (13) we shall use the estimate (8) for the current.

III. GRB ENGINE

There are three types of sites in the universe where magneic fields can induce large electric currents in the strings. They are compact structures (galaxies and clusters of galaxies), voids, and walls (filaments and sheets) of the large-scale structures. The total rate of GRBs is dominated by the walls, and further on we shall concentrate on these structures only.

Magnetic fields in our scenario are assumed to be generated in young galaxies during the bright phase of their evolution [29] and then dispersed by galactic winds in the intergalactic space. Then at present the fields are concentrated in the filaments and sheets of the large-scale structure [30,31].

Assuming that magnetic fields were generated at some $z \sim z_B$ (galaxy formation epoch) and then remained frozen in the extragalactic plasma, we obtain

$$B(z) = B_0(1+z)^2, \tag{15}$$

where the characteristic field strength at the present time B_0 can be estimated from the equipartition condition as $B_0 \sim 10^{-7}$ G [30].

With sheets of characteristic size $L \sim (20-50)h^{-1}$ Mpc and thickness $D \sim 5h^{-1}$ Mpc, we can estimate the fraction of the space occupied by the walls with magnetized plasma as $f_B \sim D/L \sim 0.1$. For numerical estimates below we shall use $z_B \sim 4$.

We shall now estimate the physical quantities characterizing GRBs powered by cusps of superconducting strings. In what follows we assume that the universe is spatially flat, is dominated by non-relativistic matter, and has age $t_0 = 0.87 \cdot 10^{10}$ yr, which corresponds to dimensionless Hubble constant $h = 0.75$.

A. GRB rate and fluence

Due to the large current, the cusp produces a powerful pulse of electromagnetic radiation. The total energy of the pulse is given by [7,8] $\mathcal{E}_{em}^{tot} \sim 2k_{em}J_0J_{max}l$, where $l \sim \alpha t$ is the length of the loop, and the coefficient $k_{em} \sim 10$ is taken from numerical calculations [7]. This radiation is emitted within a very narrow cone of openening angle $\theta_{min} \sim 1/\gamma_{max}$. The angular distribution of radiated energy at larger angles is given by [7]

$$d\mathcal{E}_{em}/d\Omega \sim k_{em}J_0^2 l/\theta^3. \tag{16}$$

For a GRB originating at redshift z and seen at angle θ with respect to the string velocity at the cusp, we have from Eqs.(8)-(15)

$$d\mathcal{E}_{em}/d\Omega \sim k_{em}e^4\alpha^3 t_0^3 B_0^2 (1+z)^{-1/2}\theta^{-3}, \tag{17}$$

The Lorentz factor of the relevant string segment near the cusp is $\gamma \sim 1/\theta$. The duration of the cusp event as seen by a distant observer is [9]

$$\tau_c \sim (1+z)(\alpha t/2)\gamma^{-3} \sim (\alpha t_0/2)(1+z)^{-1/2}\theta^3. \tag{18}$$

One can expect that the observed duration of GRB is $\tau_{GRB} \sim \tau_c$. This expectation will be justified by the hydrodynamical analysis in Section IV.

The fluence, defined as the total energy per unit area of the detector, is [10]

$$S \sim (1+z)(d\mathcal{E}_{em}/d\Omega)d_L^{-2}(z), \tag{19}$$

where $d_L(z) = 3t_0(1+z)^{1/2}[(1+z)^{1/2} - 1]$ is the luminosity distance.

The rate of GRBs originating at cusps in the redshift interval dz and seen at an angle θ in the interval $d\theta$ is given by

$$d\dot{N}_{GRB} \sim f_B \cdot \frac{1}{2}\theta d\theta (1+z)^{-1}\nu(z)dV(z). \tag{20}$$

Here, $\nu(t) \sim n_l(t)/T_l \sim 2\alpha^{-2}t^{-4}$ is the number of cusp events per unit spacetime volume, $T_l \sim \alpha t/2$ is the oscillation period of a loop, $dV = 54\pi t_0^3[(1+z)^{1/2} - 1]^2(1+z)^{-11/2}dz$ is the proper volume between redshifts z and $z+dz$, and we have used the relation $dt_0 = (1+z)dt$.

Since different cusp events originate at different redshifts and are seen at different angles, our model automatically gives a distribution of durations and fluences of GRBs. The angle θ is related to the Lorentz factor of the relevant portion of the string as $\theta \sim 1/\gamma$, and from Eqs.(17),(19) we have

$$\gamma(z;S) \sim \gamma_0 \alpha_{-8}^{-1} S_{-8}^{1/3} B_{-7}^{-2/3} [(\sqrt{1+z} - 1)^2 \sqrt{1+z}]^{1/3}. \tag{21}$$

Here, $\gamma_0 \approx 190$, $\alpha_{-8} = \alpha/10^{-8}$, and the fluence S and the magnetic field B_0 are expressed as $S = S_{-8} \cdot 10^{-8}\ erg/cm^2$ and $B_0 = B_{-7} \cdot 10^{-7}\ G$.

Very large values of $\gamma \sim \gamma_{max}$, which correspond (for a given redshift) to largest fluences, may not be seen at all because the radiation is emitted into a too narrow solid angle and the observed rates of these events are too small. The minimum value $\gamma(z;S_{min})$ is determined by the smallest fluence that is observed, *e.g.* for GRBs at $z \gtrsim 1$ with $S_{min} \sim 3 \cdot 10^{-8}$ erg/cm^2, $\gamma_{min} \approx 170$. Another lower limit on γ, which dominates at small z, follows from the condition of compactness [1] and is given by $\gamma \gtrsim 100$ (see Section IV).

The total rate of GRBs with fluence larger than S is obtained by integrating Eq.(20) over θ from $\gamma_{max}^{-1}(z)$ to $\gamma^{-1}(z;S)$ and over z from 0 to $\min[z_m; z_B]$, with z_m from $\gamma_{max}(z_m) =$

$\gamma(z_m;S)$. For relatively small fluences, $S_{-8} < S_c = 0.03(\gamma_{max}(0)\alpha_{-8}/\gamma_0)^3 B_{-7}^2$, $z_B < z_m$ and we obtain

$$\dot{N}_{GRB}(>S) \sim \frac{f_B}{2\alpha^2 t_0^4}\int_0^{z_B} dV(z)(1+z)^5\gamma^{-2}(z;S)$$
$$\sim 3\cdot 10^2 S_{-8}^{-2/3} B_{-7}^{4/3}\ yr^{-1}. \tag{22}$$

Remarkably, this rate in our model does not depend on any string parameters and is determined (for a given value of S) almost entirely by the magnetic field B_0. It agrees with the observed rate for $B_{-7} \sim 1$ (formally, the observed rate $\dot{N} \sim 300\text{yr}^{-1}$ at $S > 1\cdot 10^{-7}\text{erg/cm}^2$ gives $B_{-7} = 3.2$). The predicted slope $\dot{N}_{GRB}(>S) \propto S^{-2/3}$ is also in a reasonable agreement with the observed one $\dot{N}_{obs}(>S) \propto S^{-0.55}$ at relatively small fluences [32].

For large fluences $S_{-8} > S_c$, integration of Eq.(20) gives $\dot{N}_{GRB}(>S) \propto S^{-3/2}$. Observationally, the transition to this regime occurs at $S_{-8} \sim 10^2 - 10^3$. This can be accounted for if the cusp development is terminated by small-scale wiggles with fractional energy in the wiggles $\epsilon \sim 10^{-7}\alpha_{-8}^2 B_{-7}^{4/3}$. Alternatively, if γ_{max} is determined by the back-reaction of the charge carriers, Eq.(10), then the regime (22) holds for larger S_{-8}, and observed steepening of the distribution at large S can be due to the reduced efficiency of BATSE to detection of bursts with large γ. Indeed, large γ results in a large Lorentz factor γ_{CD} of the emitting region (see Section IV), and at $\gamma_{CD} \gtrsim 10^3$ photons start to escape from the BATSE range.

B. GRB duration

The duration of GRBs originating at redshift z and having fluence S is readily calculated from Eqs.(18) and (21) as

$$\tau_{GRB} \approx 200\frac{\alpha_{-8}^4 B_{-7}^2}{S_{-8}}(1+z)^{-1}(\sqrt{1+z}-1)^{-2}\ s \tag{23}$$

From Eqs.(20) and (18) we find the rate of GRBs with durations in the interval $d\tau$ and redshifts in the interval dz,

$$d\dot{N} \sim 10^2 \alpha^{-2} t_0^{-1} \left(\frac{\tau}{\alpha t_0}\right)^{2/3} (\sqrt{1+z}-1)^2 (1+z)^{-1/6} dz \frac{d\tau}{\tau}. \tag{24}$$

The distribution of GRB durations is found by integrating this over z. The integration is restricted by $z < z_B$ and $S > S_{min} \sim 3 \cdot 10^{-8}$ erg/cm^2. The latter condition can be expressed as $z < \tilde{z}(\tau)$, where $\tilde{z}(\tau)$ is the solution of Eq. (23) for z with $S \sim S_{min}$.

The distribution changes its form at the characteristic value τ_* defined by $\tilde{z}(\tau_*) = z_B$. With $z_B = 4$, Eq. (23) gives

$$\tau_* \sim 8.7 \alpha_{-8}^4 B_{-7}^2 \text{ s}. \tag{25}$$

For $\tau < \tau_*$ we have $d\dot{N} \propto \tau^{2/3} d\tau/\tau$, and for $\tau \gg \tau_*$, $\quad d\dot{N} \propto \tau^{-5/6} d\tau/\tau$. We thus see that the distribution is peaked at $\tau \sim \tau_*$.

The largest value of τ in our model is obtained from Eq.(18) with $\theta \sim \theta_{max} \sim 10^{-2}$, $\tau_{GRB}^{max} \sim 10^3 \alpha_{-8}$ s. There is no sharp lower cutoff, but very small values of τ will not be observed due to the low rate of events. With the rate $\sim 10^2$ yr^{-1} near the peak of the distribution, the rate of events with $\tau \sim 10^{-4}\tau_*$ is about 0.1 yr^{-1}.

A lower bound on τ is also set by the detector resolution ($\sim 10^{-2}$ s for BATSE). Hence, we have $\tau_{min} \sim \max\{10^{-4}\tau_*, 10^{-2} \text{ s}\}$.

The observed distribution of GRB durations extends from $\sim 10^{-2}$ s to $\sim 10^3$ s. The distribution is bimodal, with peaks at 0.5 s and 15 s [33], and there are some observational indications that short and long GRBs may have different origin. Our model is probably better suited to describe the short GRB population (see Section V). With $\tau_* \sim 0.5$ s and $B_{-7} \sim 3$, Eq.(25) gives $\alpha_{-8} \sim 0.3$. This corresponds to the string symmetry breaking scale $\eta \sim 1 \cdot 10^{14}$ GeV. The range of GRB durations is then given by $\tau_{GRB}^{min} \sim 10^{-2}$ s, $\tau_{GRB}^{max} \sim 10^3$ s.

It should be noted that the validity of our simple one-parameter model does not extend beyond rough order-of-magnitude estimates (see Section VI). In particular, it is not expected to give the correct duration distribution $\dot{N}(\tau)$, and identifying the peaks of the theoretical and observed distributions may therefore exceed the accuracy of the model. A more con-

servative approach is to require that the characteristic duration τ_* lies within the observed range of GRB durations. This gives $0.2 \lesssim \alpha_{-8} \lesssim 3$.

IV. ACCELERATION AND HYDRODYNAMICS

A beam of low-frequency e-m radiation propagating in plasma produces a beam of accelerated particles.

The characteristic frequency of e-m radiation in a pulse produced by a cusp segment with Lorentz factor γ is

$$\omega_{em} \sim \frac{4\pi}{\alpha t_0}\gamma^3(1+z)^{3/2} = 4.6(\gamma/10^3)^3(1+z)^{3/2}\alpha_{-8}^{-1}\ \mathrm{s}^{-1}. \tag{26}$$

The plasma frequency in the intergalactic gas of density $n = n_{-5}10^{-5}\ \mathrm{cm}^{-3}$,

$$\omega_{pl} = 1.8 \cdot 10^2 n_{-5}^{1/2}\ \mathrm{s}^{-1}, \tag{27}$$

is higher than ω_{em} when $\gamma \lesssim 3.4 \cdot 10^3 n_{-5}^{1/6}\alpha_{-8}^{1/3}(1+z)^{-1/2}$. Therefore, the low-frequency radiation from the cusp cannot propagate in plasma. In fact, the energy density of e-m beam is much larger than that of the plasma, and the beam would push the plasma away even in the case $\omega_{em} > \omega_{pl}$. This process occurs due to the acceleration of plasma particles.

Let us consider the propagation of a charged test particle in a strong, low-frequency e-m wave. For a time interval much shorter than the period of the wave, $t \ll 1/\omega_{em}$, the e-m field of the wave can be approximated by static, orthogonal electric and magnetic fields of equal magnitude. Solution of the equations of motion (see e.g. [34]) shows that both positive and negative charges are accelerated mainly in the direction of wave propagation, $\mathbf{n} = (\mathbf{E} \times \mathbf{B})/EB$, with their Lorentz factor increasing with time as

$$\gamma_b(t) = \left(\frac{3}{\sqrt{2}}\frac{eB}{m}t\right)^{2/3}, \tag{28}$$

where m is the particle's mass. The synchrotron energy loss of an accelerated particle is small, because when it moves in the direction of wave propagation, $\mathbf{n}$, the electric force, $e\mathbf{E}$,

and the magnetic force $e\mathbf{v} \times \mathbf{B}$ almost exactly compensate each other: $e(\mathbf{E} + v\mathbf{n} \times \mathbf{B}) \approx 0$. This regime of acceleration is practically the same as in the Gunn-Ostriker mechanism [35].

For an e-m wave in vacuum, a test particle would be accelerated at $t \sim 1/\omega_{em}$ up to a very large Lorentz factor. But the maximum Lorentz factor of the *beam* is saturated at the value γ_b, when the energy of the beam reaches the energy of the original e-m pulse: $N_b m \gamma_b \sim \mathcal{E}_{em}$. This results in the Lorentz factor of the beam

$$\gamma_b \sim 4 \cdot 10^2 B_{-7}^2 n_{-5}^{-1} (1+z)^4 (\gamma/100)^6. \tag{29}$$

Let us now turn to the hydrodynamical phenomena in which the gamma radiation of the burst is actually generated. The beam of accelerated particles pushes the gas with the frozen magnetic field ahead of it, producing an external shock in surrounding plasma and a reverse shock in the beam material, as in the case of "ordinary" fireball (for a review see [1]). The difference is that the beam propagates with a very large Lorentz factor $\gamma_b \gg \gamma$, where γ is the Lorentz factor of the cusp (the precise value of γ_b is not important for our discussion). Another difference is that the beam propagates in a very low-density gas. The beam can be regarded as a narrow shell of relativistic particles of width $\Delta \sim l/2\gamma^3$ in the observer's frame.

The gamma radiation of the burst is produced as synchrotron radiation of electrons accelerated by external and reverse shocks. Naively, the duration of synchrotron radiation, i.e. τ_{GRB}, is determined by the thickness of the shell as $\tau_{GRB} \sim \Delta$. This is confirmed by a more detailed analysis, as follows. The reverse shock in our case is ultrarelativistic [36,1]. The neccessary condition for that, $\rho_b/\rho < \gamma_b^2$, is satisfied with a wide margin (here ρ_b is the baryon density in the beam and ρ is the density of unperturbed gas). In this case, the shock dynamics and the GRB duration are determined by two hydrodynamical parameters [1]. They are the thickness of the shell Δ and the Sedov length, defined as the distance travelled by the shell when the mass of the snow-ploughed gas becomes comparable to the initial energy of the beam. The latter is given by $l_{Sed} \sim (\mathcal{E}_{iso}/\rho)^{1/3}$.

The reverse shock enters the shell and, as it propagates there, it strongly decelerates the shell. The synchrotron radiation occurs mainly in the shocked regions of the shell and of the external plasma. The surface separating these two regions, the contact discontinuity (CD) surface, propagates with the same velocity as the shocked plasma, where the GRB radiation is produced.

The synchrotron radiation ceases when the reverse shock reaches the inner boundary of the shell. This occurs at a distance $R_\Delta \sim l_{Sed}^{3/4}\Delta^{1/4}$ when the Lorentz factor of the CD surface is

$$\gamma_{CD} \sim (l_{Sed}/\Delta)^{3/8} \sim 0.1 B_{-7}^{1/4} n_{-5}^{-1/8} (1+z)^{1/2} \gamma^{3/2}. \tag{30}$$

Note that these values do not depend on the Lorentz factor of the beam γ_b and are determined by the cusp Lorentz factor γ. The size of the synchrotron emitting region is of the order R_Δ, and the Lorentz factor of this region is equal to γ_{CD}. The compactness condition [1] requires $\gamma_{CD} \gtrsim 10^2$, and Eq. (30) yields $\gamma \gtrsim 10^2$ which we used earlier in Section III. The duration of GRB is given by

$$\tau_{GRB} \sim R_\Delta/2\gamma_{CD}^2 \sim l/2\gamma^3, \tag{31}$$

i.e. it is equal to the duration of the cusp event given by Eq.(18). The energy that goes into synchrotron radiation is comparable to the energy of the electromagnetic pulse.

V. PREDICTIONS AND PROBLEMS

In this section we shall consider a number of predictions of our model. Some of these predictions pose potential problems.

(i) Short-time structure of GRBs.

Most of GRBs exhibit a complex short-time structure. These variations must be a property of inner engine [1,37]. In the cusp model they can be naturally produced by wiggles.

Wiggles are amplified in near-cusp regions and, acting like minicusps, produce a sequence of successive fireballs. A quantitative analysis of this effect would require a detailed study of the gravitational back-reaction, which controls the amplitude of the wiggles.

(ii) Repeaters.

Cusps reappear on a loop with a period of loop oscillation, producing nearly identical GRBs. In our model, where all loops have the same length $l = \alpha t$ at a given cosmological epoch t, the recurrence time, $T_l \sim (1+z)\alpha t/2 \sim 50\alpha_{-8}(1+z)^{-1/2}$ yr, is too long to be observed by BATSE and other detectors. In a more realistic models, some fraction of loops would have lengths smaller than αt and thus shorter recurrence periods. This fraction is model-dependent. Moreover, GRBs from repeaters with $l < \alpha t$ must be weak and have short durations.

The GRB fluence from a loop of length l produced by a string segment with Lorentz factor γ can be readily calculated as

$$S \sim \frac{1}{9} k_{em} e^4 l^3 \gamma^3 \frac{B_0^2}{t_0^2} \frac{(1+z)^4}{[(1+z)^{1/2}-1]^2} \tag{32}$$

After a change of variables from γ and l to τ_{GRB} and the reccurence period $T_{rec} = l(1+z)/2$, we obtain

$$\tau_{GRB} \sim k_{em} e^4 \frac{B_0^2}{t_0^2} \frac{(1+z)}{[(1+z)^{1/2}-1]^2} \frac{T_{rec}^4}{S}. \tag{33}$$

A search for repeaters with $T_{rec} \leq 5$ yr requires, according Eq.(33), low fluences $S \lesssim 10^{-7}$ erg/cm^2 and short durations $\tau_{GRB} \lesssim 40$ ms. The BATSE efficiency is low for such events [38] and the repeating burst could easily have been lost. The total number of GRBs shorter than 40 ms in BATSE 3B catalogue is less than 5 [39].

(iii) Host galaxies

The discovery of GRB aferglows revealed an association of long-duration GRBs with galaxies (see [40] for a review). 19 GRBs with long durations are found to be undoubtedly

hosted by normal galaxies [41]. For 10 of them redshifts are found to be typically 1 - 3. For many bright bursts, which are most probably at small distances, no host galaxies have been found. For example, for 16 bright bursts observed by the Interplanetary Network with small error boxes, no galaxies are found with magnitudes from 20 to 24. This suggests that some of the GRBs are not hosted by galaxies.

In our model, the fraction of loops captured by galaxes is expected to be small, due to the high velocities of the loops. The most straightforward way to reconcile the model with observations is to assume that cusps are responsible only for a subset of the observed GRBs not associated with galaxies. Such a subset could include the short-duration GRBs, for which no host galaxies have yet been detected. With the choice of parameters $B_{-7} \sim 3$, $\alpha_{-8} \sim 0.3$, as in Section III, the distribution of GRB durations is peaked at $\tau_{GRB} \sim 0.5$ s. At the same time, the tail of the distribution extends all the way to $\tau_{GRB} \sim 10^3$ s, and thus the model can account for *some* of the long GRBs as well. This particular possibility meets another problem, since short GRBs do not show deviation from Euclidean distribution. However, it is often suggested that GRBs comprise a few subclasses, and the existence of a no-host subclass remains plausible.

An alternative possibility should be also mentined. In string evolution models with $\alpha \gg k_g G\mu$, the lifetime of the loops is $\tau_l \gg t$, so the loops will be slowed down by the expansion of the universe and a substantial fraction of them can be captured in galaxies.

(iv) Bursts of gravitational radiation

Our model predicts that GRBs should be accompanied by strong bursts of gravitational radiation (see also [42]). The angular distribution of the gravitational wave energy around the direction of the cusp is [43] $d\mathcal{E}_g/d\Omega \sim G\mu^2 l/\theta$, and the dimensionless amplitude of a burst of duration τ originating at redshift z can be estimated as

$$h \sim k_g^{-1/2} \alpha^{5/3} (\tau/t_0)^{1/3} (1+z)^{-1/3} z^{-1}, \tag{34}$$

or $h \sim 10^{-21} \alpha_{-8}^{5/3} z^{-1} (\tau/1s)^{1/3}$ for $z \lesssim 1$. Here, we have used the relation $F_g \sim h^2/G\tau^2 \sim$

$(1+z)(d\mathcal{E}_g/d\Omega)/d_L^2\tau$ for the gravitational wave flux and Eq.(18) for the burst duration τ. These gravitational wave bursts are much stronger than expected from more conventional sources and should be detectable by the planned LIGO, VIRGO and LISA detectors. It has been shown in [42] that gravitational wave bursts from strings are linearly polarized and have a characteristic waveform $h(t) \propto t^{1/3}$.

(v) X- and γ- ray diffuse radiation

Tremendous energy [see Eq.(16)] released in a narrow angle $\theta \sim 1/\gamma_{max}$ is not seen in GRBs because of the smallness of this angle. The beam of particles accelerated by e-m radiation in this narrow cone has a very large Lorentz factor, and the emitted photons have energies in excess of 1 TeV. These photons are absorbed in collisions with infrared (IR) or microwave photons, collectively denoted as γ_t: $\gamma + \gamma_t \to e^+ + e^-$. Electrons and positrons start e-m cascades on microwave photons (γ_{bb}) due to Inverse Compton scattering ($e+\gamma_{bb} \to e+\gamma$) and pair production ($\gamma+\gamma_{bb} \to e^+ + e^-$). As they degrade in energy, cascade electrons are effectively deflected in the extragalactic magnetic field, and the produced diffuse gamma radiation is isotropic. The spectrum of remaining cascade photons was calculated analytically in [44] (for recent Monte Carlo simulation see [45]). The analytic spectrum is described in terms of three parameters: ϵ_γ, ϵ_X and ω_γ.

ϵ_γ is the minimum energy of absorption, *i.e.* the smallest energy of a photon absorbed on IR radiation ($\epsilon_\gamma \sim m_e^2/\varepsilon_{IR}$, with the exact value dependent on the spectrum of IR radiation). ϵ_X is the energy of an IC photon produced by an electron of energy $\epsilon_e = \epsilon_\gamma/2$, *i.e.* by an electron born by a photon of energy ϵ_γ. ω_γ is the energy density of cascade radiation.

The space density of cascade photons, $n_\gamma(E)$, is given by [44]

$$n_\gamma(E) = \begin{cases} K(E/\epsilon_X)^{-3/2} & \text{if} \quad E \le \epsilon_X \\ K(E/\epsilon_X)^{-2} & \text{if} \quad \epsilon_X \le E \le \epsilon_\gamma \\ 0 & \text{if} \quad E > \epsilon_\gamma \end{cases} \tag{35}$$

where K is a normalization constant which can be expressed in terms of ω_{cas} as

$$K = \frac{\omega_{cas}}{\epsilon_X^2 (2 + \ln \epsilon_\gamma / \epsilon_X)} \tag{36}$$

The cascade energy density ω_{cas} can be calculated as the total energy release in e-m radiation of cusps integrated over redshits from 0 up to $z_B = 4$. This gives

$$\omega_{cas} = 5 f_B k_{em} e^3 \eta B_0 / t_0 \tag{37}$$

Assuming $\epsilon_\gamma \approx 100$ GeV due to absorption on IR radiation, we obtain $\epsilon_X \approx 8.1$ MeV. Then the predicted spectrum in the energy range 8 MeV $\leq E \leq$ 100 GeV is

$$I_{theor}(E) \sim 2.5 \cdot 10^{-10} \alpha_{-8}^{1/2} B_{-7} \left(\frac{E}{10^3 \mathrm{MeV}} \right)^{-2} \; cm^{-2} s^{-1} sr^{-1} MeV^{-1}. \tag{38}$$

This is to be compared with the EGRET flux [46] for the energy range 5 MeV $\leq E \leq$ 100 GeV,

$$I_{obs}(E) = 1.38 \cdot 10^{-9} \left(\frac{E}{10^3 \ \mathrm{MeV}} \right)^{-2.1 \pm 0.03} \; cm^{-2} s^{-1} sr^{-1} MeV^{-1}. \tag{39}$$

With $\alpha_{-8} \sim 0.3$ and $B_{-7} \sim 3$ the predicted flux differs from the observed one by a factor of 3. This can be regarded as agreement for an order of magnitude estimate of our simple model.

(vi) Ultra High Energy Cosmic Rays

GRBs have been suggested as possible sources of the observed ultrahigh-energy cosmic rays (UHECR) [47,48]. This idea encounters two difficulties. First, if GRBs are distributed uniformly in the universe, UHECR have a classical Greisen-Zatsepin-Kuzmin (GZK) cutoff, absent in the observations. Second, the acceleration by an ultrarelativistic shock is possible only in the one-loop regime (*i.e.* due to a single reflection from the shock) [49]. For a standard GRB with a Lorentz factor $\gamma_{sh} \sim 300$ it results in the maximum energy $E_{max} \sim \gamma_{sh}^2 m_p \sim 10^{14}\ eV$, far too low for UHECR.

Our model can resolve both of these difficulties, assuming that γ_{max} is determined by the current backreaction, Eq.(10).

If the magnetic field in the Local Supercluster (LS) is considerably stronger than outside, then the cusps in LS are more powerful and the GZK cutoff is less pronounced.

Cusp segments with large Lorentz factors produce hydrodinamical flows with large Lorentz factors, *e.g.*, $\gamma \sim 2{\cdot}10^4$ corresponds to $\gamma_{CD} \sim 3{\cdot}10^5$ and $E_{max} \sim \gamma_{CD}^2 m_p \sim 1{\cdot}10^{20}\ eV$. Protons with such energies are deflected in the magnetic field of LS and can be observed, while protons with much higher energies caused by near-cusp segments with $\gamma \gtrsim 10^5$ are propagating rectilinearly and generally are not seen. A quantitative analysis of the UHECR flux in this scenario will be given elsewhere.

VI. DISCUSSION AND CONCLUSIONS

The nature of GRB engines is still unknown. There are observational indications that they are astrophysical objects: about 25 GRBs are reliably found to be located in galaxies, probably in regions of star formation; at least in one case GRB is identified with a supernova (SN 1998 bw). The most popular now are astrophysical models, with binary neutron star mergers [50], failed supernova [51], hypernova [52] and supranova [53] being the front runners (for a critical review see [54]). All these models, however, are not developed enough to give quantitative predictions. They also share the difficulty with explaining the large beaming factor required for GRBs.

In contrast, the cosmic string model presented here allows one to obtain quantitative predictions for the main observational characteristics of GRBs. In this paper we developed a deliberately simplified model, which is characterized by a single free parameter (the energy scale η of symmetry breaking, or $\alpha = k_g G\eta^2$) and by three other physical quantities, relatively well restricted (the magnetic field in filaments and sheets B_0, the epoch of galaxy formation z_B, and the density of baryonic matter in the filaments and sheets, a quantity not critical for the predictions). Nevertheless, the model correctly accounts for the GRB rate, and for the range of GRB fluences and durations. It may also explain the short-time structure of

GRBs, the diffuse X- and γ-ray backgrounds, and the ultrahigh-energy cosmic rays.

The string model predicts recurrence of GRBs with a period of $T_l \sim 50\alpha_{-8}(1+z)^{-1/2}$ yrs. Very short bursts may have much shorter recurrence periods, perhaps as short as a few years. Observation of these repeaters is a challenge for the future detectors with a high efficiency for detection of short bursts.

Another testable prediction of the model is that GRBs should be accompanied by strong bursts of gravitational radiation with a characteristic waveform.

It must be emphasized that our model involves a number of simplifying assumptions. All loops at cosmic time t were assumed to have the same length $l \sim \alpha t$ with $\alpha \sim k_g G\eta^2$, while in reality there should be a distribution $n(l,t)$. The evolution law (15) for $B(z)$ and the assumption of $f_B =$ const are also oversimplified. A more realistic model should also account for a spatial variation of B. Being basically a one-parameter model, our model may predict spurious correlations between the GRB characteristics. In particular, the $S \propto \tau_{GRB}^{-1}$ correlation, suggested by Eq.(23), holds only at a fixed redshift and tends to be washed out when the redshit distribution, the loop length distribution $n(l,t)$, and the inhomogeneous spatial distribution of the magnetic field are taken into account.

Our model meets basically one difficulty: it predicts too low GRB rate from galaxies. This discrepancy could be explained if our model strongly underestimates the capture rate of string loops by galaxies. For example, if $\alpha \gg k_g G\mu$, then the loops are non-relativistic and may be effectively captured by galaxies. Another possibility is that our model could describe some subclass of the sources not associated with galaxies. Such a subclass could include the short-duration GRBs for which host galaxies are not found. In this case, the model needs a smaller α, as discussed in section IIIB. In contrast to the prediction of our model, the short bursts do not show strong deviation from the Euclidean distribution. This could be due to observational selection effects, since the faint short-duration GRBs which form this subclass have a low detection efficiency in BATSE. Alternatively, it could be another subclass of

no-host GRBs.

On the other hand, GRBs from cusps have properties which distinctly distinguish them from those produced by collapsars: they are periodically repeating on the scale of a few decades for majority of GRBs and on the scale of a few years for faint bursts ($S \lesssim 10^{-8}$ erg/cm^2) with short duration $\tau_{\rm GRB} \lesssim 20$ ms. The next generation of GRB detectors can examine this prediction.

ACKNOWLEDGEMENT

We are grateful to Roger Blandford, Ken Olum and Bohdan Paczynski for useful discussions. The work of VB and BH was supported in part by INTAS through grant No 1065 and the work of AV by the National Science Foundation (USA).

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